

# TeV-Scale Seesaw with Loop-Induced Dirac Mass Term and Dark Matter from $U(1)_{B-L}$ Gauge Symmetry Breaking

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## Abstract

We show a TeV-scale seesaw model where Majorana neutrino masses, the dark matter mass, and stability of the dark matter can be all originated from the  $U(1)_{B-L}$  gauge symmetry. Dirac mass terms for neutrinos are forbidden at the tree level by  $U(1)_{B-L}$ , and they are induced at the one-loop level by spontaneous  $U(1)_{B-L}$  breaking. The right-handed neutrinos can be naturally at the TeV-scale or below because of the induced Dirac mass terms with loop suppression. Such right-handed neutrinos would be discovered at the CERN Large Hadron Collider (LHC). On the other hand, stability of the dark matter is guaranteed without introducing an additional  $Z_2$  symmetry by a remaining global  $U(1)$  symmetry after the  $U(1)_{B-L}$  breaking. A Dirac fermion  $\Psi_1$  or a complex neutral scalar  $s_1^0$  is the dark matter candidate in this model. Since the dark matter ( $\Psi_1$  or  $s_1^0$ ) has its own B–L charge, the invisible decay of the  $U(1)_{B-L}$  gauge boson  $Z'$  is enhanced. Experimental constraints on the model are considered, and the collider phenomenology at the LHC as well as future linear colliders is discussed briefly.

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## I. INTRODUCTION

Neutrino oscillation measurements [1–5] have established evidence for tiny neutrino masses, which are supposed to be zero in the standard model (SM) of particle physics. It seems mysterious that the scale of neutrino masses is much smaller than that of the other fermion masses. The simplest way to obtain tiny neutrino masses is the seesaw mechanism [6] where right-handed neutrinos are introduced. Due to suppression with huge Majorana masses of the right-handed neutrinos as compared to the electroweak scale, neutrino masses can be very small even though Dirac Yukawa coupling constants for neutrinos are of  $\mathcal{O}(1)$ . However testability of the mechanism seems to be a problem because key particles (right-handed neutrinos) with such huge masses would not be accessible in future experiments. A possible solution for the problem is radiative generation of Dirac Yukawa couplings for neutrinos. In order to explain tiny neutrino masses, the right-handed neutrinos with masses of  $\mathcal{O}(100)$  GeV are acceptable naturally without assuming excessive fine tuning among coupling constants by virtue of loop-suppressed Dirac Yukawa couplings. Radiative generation of Dirac Yukawa couplings has been discussed in various frameworks such as the left-right symmetry [8, 9], supersymmetry (SUSY) [10], extended models within the SM gauge group [11–13], and an extra  $U(1)$  gauge symmetry [14] (See also ref. [15]).

On the other hand, existence of dark matter has been indicated by Zwicky [7], and its thermal relic abundance has been quantitatively determined by the WMAP experiment [16]. If the essence of the dark matter is an elementary particle, the weakly interacting massive particle (WIMP) would be a promising candidate. A naive power counting shows that the WIMP dark matter mass should be at the electroweak scale. This would suggest a strong connection between the WIMP dark matter and the Higgs sector. It is desired to have viable candidate for dark matter in models beyond the standard model. Usually stability of the dark matter candidate is ensured by imposing a  $Z_2$  parity where the dark matter is the lightest  $Z_2$ -odd particle. It is well known that such  $Z_2$  odd particles are compatible with radiative neutrino mass models [14, 17–21]. Usually in such models, however, the origin of the  $Z_2$  parity has not been clearly discussed. It seems better if a global symmetry to stabilize the dark matter is not just imposed additionally but obtained as a remnant of some broken symmetry which is used also for other purposes [19].

In this paper we propose a new model in which tiny neutrino masses and the origin of

Particles	$s^0$	$\eta$	$(\Psi_R)_i$	$(\Psi_L)_i$	$(\nu_R)_i$	$\sigma^0$
SU(3) <sub>C</sub>	<u>1</u>	<u>1</u>	<u>1</u>	<u>1</u>	<u>1</u>	<u>1</u>
SU(2) <sub>L</sub>	<u>1</u>	<u>2</u>	<u>1</u>	<u>1</u>	<u>1</u>	<u>1</u>
U(1) <sub>Y</sub>	0	1/2	0	0	0	0
U(1) <sub>B-L</sub>	1/2	1/2	-1/2	3/2	1	2

TABLE I: New particles and their properties under gauge symmetries of the model.

dark matter are naturally explained at the TeV-scale. We introduce the  $U(1)_{B-L}$  gauge symmetry to the SM gauge group which is spontaneously broken at multi-TeV scale [22–24]. Its collider phenomenology has been studied [25–28]. In our model, Dirac Yukawa couplings for neutrinos are forbidden at the tree level by the  $U(1)_{B-L}$ . They are generated at the one-loop level after the  $U(1)_{B-L}$  breaking. Simultaneously, right-handed neutrino masses are generated at the tree level by spontaneous breaking of the  $U(1)_{B-L}$ . As a result, light neutrino masses are obtained effectively at the two-loop level without requiring too small coupling constants ( $\lesssim 10^{-3}$ ) from the TeV-scale physics. Furthermore it turns out that the dark matter in our model is stabilized by an unbroken global  $U(1)$  symmetry which appears automatically in the Lagrangian with appropriate assignments of the  $U(1)_{B-L}$ -charges for new particles. The mass of a dark matter candidate  $\Psi_1$ , which is a Dirac fermion, is also generated by the  $U(1)_{B-L}$  breaking (See also ref. [29]<sup>1</sup>). We show that the model is viable under the current experimental constraints. Prospects in collider experiments are also discussed.

## II. THE MODEL

In our model, the  $U(1)_{B-L}$  gauge symmetry is added to the SM gauge group. New particles and their properties under gauge symmetries of the model are shown in Table I. Fields  $s^0$ ,  $\eta$ , and  $\sigma^0$  are complex scalars while  $(\Psi_R)_i$ ,  $(\Psi_L)_i$ , and  $(\nu_R)_i$  ( $i = 1, 2$ ) are Weyl fermions. All of them except  $\eta [= (\eta^+, \eta^0)^T]$  are singlet fields under the SM gauge group. The SM Higgs doublet field  $\Phi [= (\phi^+, \phi^0)^T]$  and  $\eta$  have different  $U(1)_{B-L}$ -charges although

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<sup>1</sup> The dark matter in ref. [29] does not contribute to the mechanism to generate light neutrino masses.

their representations for the SM gauge group are the same. Notice that mass terms of  $\Psi_R$ ,  $\Psi_L$ , and  $\nu_R$  are forbidden by the  $U(1)_{B-L}$  symmetry.

Yukawa interactions are given by

$$\begin{aligned} \mathcal{L}_{\text{Yukawa}} = & \mathcal{L}_{\text{SM-Yukawa}} - (y_R)_i \overline{(\nu_R)_i^c} (\nu_R)_i (\sigma^0)^* - (y_\Psi)_i \overline{(\Psi_R)_i} (\Psi_L)_i (\sigma^0)^* \\ & - (y_3)_{ij} \overline{(\nu_R)_i^c} (\Psi_R)_j (s^0)^* - h_{ij} \overline{(\Psi_L)_i} (\nu_R)_j s^0 - f_{\ell i} \overline{(L_L)_\ell} (\Psi_R)_i i\sigma_2 \eta^* + \text{h.c.}, \end{aligned} \quad (1)$$

where  $\mathcal{L}_{\text{SM-Yukawa}}$  stands for the Yukawa interactions in the SM and  $(L_L)_\ell$  are the lepton doublet fields of flavor  $\ell$  ( $\ell = e, \mu, \tau$ ). Superscript  $c$  means the charge conjugation and  $\sigma_i$  ( $i = 1-3$ ) are the Pauli matrices. We take a basis where Yukawa matrices  $y_R$  and  $y_\Psi$  are diagonalized such that their real positive eigenvalues satisfy  $(y_R)_1 \leq (y_R)_2$  and  $(y_\Psi)_1 \leq (y_\Psi)_2$ .

Scalar potential in this model is expressed as

$$\begin{aligned} V(\Phi, s, \eta, \sigma) = & -\mu_\phi^2 \Phi^\dagger \Phi + \mu_s^2 |s^0|^2 + \mu_\eta^2 \eta^\dagger \eta - \mu_\sigma^2 |\sigma^0|^2 \\ & + \lambda_\phi (\Phi^\dagger \Phi)^2 + \lambda_s |s^0|^4 + \lambda_\eta (\eta^\dagger \eta)^2 + \lambda_\sigma |\sigma^0|^4 \\ & + \lambda_{s\eta} |s^0|^2 \eta^\dagger \eta + \lambda_{s\phi} |s^0|^2 \Phi^\dagger \Phi + \lambda_{\phi\phi} (\eta^\dagger \eta) (\Phi^\dagger \Phi) + \lambda_{\eta\phi} (\eta^\dagger \Phi) (\Phi^\dagger \eta) \\ & + \lambda_{s\sigma} |s^0|^2 |\sigma^0|^2 + \lambda_{\sigma\eta} |\sigma^0|^2 \eta^\dagger \eta + \lambda_{\sigma\phi} |\sigma^0|^2 \Phi^\dagger \Phi + (\mu_3 s^0 \eta^\dagger \Phi + \text{h.c.}), \end{aligned} \quad (2)$$

where  $\mu_\phi^2$ ,  $\mu_\sigma^2$ ,  $\mu_s^2$ , and  $\mu_\eta^2$  are positive values. The coupling constant  $\mu_3$  of the trilinear term can be taken as a real positive value by redefinition of phase of  $s^0$ . The  $\mu_3$  is the breaking parameter for a global  $U(1)_{\eta-L}$  which conserves difference between the  $\eta$  number and the SM lepton number. Some coupling constants in the potential are constrained by the tree-level unitarity [30]. Notice that this model has a global  $U(1)$  symmetry (we refer to it as  $U(1)_{\text{DM}}$ ) of which  $s^0$ ,  $\eta$ ,  $\Psi_R$ , and  $\Psi_L$  have the same charge and others have no charge. Because of the global  $U(1)_{\text{DM}}$  symmetry, the Lagrangian is not changed by an overall shift of  $U(1)_{B-L}$ -charges with an integer for the  $U(1)_{\text{DM}}$ -charged particles.

The  $U(1)_{B-L}$  is broken spontaneously by the vacuum expectation value (vev) of  $\sigma^0$ ,  $v_\sigma [= \sqrt{2}\langle\sigma^0\rangle]$  while the  $SU(2)_L \times U(1)_Y$  is broken to  $U(1)_{\text{EM}}$  by the vev of  $\phi^0$ ,  $v [= \sqrt{2}\langle\phi^0\rangle \simeq 246 \text{ GeV}]$ . By imposing the stationary condition,  $v_\sigma$  and  $v$  are determined as

$$\begin{pmatrix} v^2 \\ v_\sigma^2 \end{pmatrix} = \frac{1}{\lambda_\sigma \lambda_\phi - \lambda_{\sigma\phi}^2/4} \begin{pmatrix} \lambda_\sigma & -\lambda_{\sigma\phi}/2 \\ -\lambda_{\sigma\phi}/2 & \lambda_\phi \end{pmatrix} \begin{pmatrix} \mu_\phi^2 \\ \mu_\sigma^2 \end{pmatrix}. \quad (3)$$

The gauge boson  $Z'$  of  $U(1)_{B-L}$  acquires its mass as  $m_{Z'} = 2g_{B-L}v_\sigma$ , where  $g_{B-L}$  denotes gauge coupling constant of  $U(1)_{B-L}$ . A constraint  $v_\sigma > 3.5 \text{ TeV}$  is given by precision tests of

the electroweak interaction [31]. Furthermore, right-handed neutrinos  $(\nu_R)_i$  obtain Majorana masses  $(m_R)_i [= \sqrt{2}(y_R)_i v_\sigma]$  while  $(\Psi_R)_i$  and  $(\Psi_L)_i$  for each  $i$  become a Dirac fermion  $\Psi_i$  with its mass  $m_{\Psi_i} [(y_\Psi)_i v_\sigma / \sqrt{2}]$ . Since the global  $U(1)_{\text{DM}}$  is not broken by  $v_\sigma$ , the lightest  $U(1)_{\text{DM}}$ -charged particle is stable. Notice that there is no anomaly for the  $U(1)_{\text{DM}}$  because  $(\Psi_R)_i$  and  $(\Psi_L)_i$  have the same  $U(1)_{\text{DM}}$ -charge. If the particle is electrically neutral ( $\Psi_1$  or a mixture of  $s^0$  and  $\eta^0$ ), it becomes a candidate for the dark matter.

After symmetry breaking with  $v_\sigma$  and  $v$ , mass eigenstates of two CP-even scalars and their mixing angle  $\alpha$  are given by

$$\begin{pmatrix} h^0 \\ H^0 \end{pmatrix} = \begin{pmatrix} \cos \alpha & -\sin \alpha \\ \sin \alpha & \cos \alpha \end{pmatrix} \begin{pmatrix} \phi_r^0 \\ \sigma_r^0 \end{pmatrix}, \quad \sin 2\alpha = \frac{2\lambda_{\sigma\phi} v v_\sigma}{m_{H^0}^2 - m_{h^0}^2}, \quad (4)$$

where  $\sigma^0 = (v_\sigma + \sigma_r^0 + iz_\sigma)/\sqrt{2}$  and  $\phi^0 = (v + \phi_r^0 + iz_\phi)/\sqrt{2}$ . It is needless to say that  $z_\phi$  and  $z_\sigma$  are Nambu-Goldstone bosons which are absorbed by  $Z$  and  $Z'$ , respectively. Masses of  $h^0$  and  $H^0$  are defined by

$$\begin{aligned} m_{h^0}^2 &= \lambda_\phi v^2 + \lambda_\sigma v_\sigma^2 - \sqrt{(\lambda_\phi v^2 - \lambda_\sigma v_\sigma^2)^2 + \lambda_{\sigma\phi}^2 v^2 v_\sigma^2}, \\ m_{H^0}^2 &= \lambda_\phi v^2 + \lambda_\sigma v_\sigma^2 + \sqrt{(\lambda_\phi v^2 - \lambda_\sigma v_\sigma^2)^2 + \lambda_{\sigma\phi}^2 v^2 v_\sigma^2}. \end{aligned} \quad (5)$$

On the other hand, since  $s^0$  and  $\eta^0$  are  $U(1)_{\text{DM}}$ -charged particles, they are not mixed with  $\sigma^0$  and  $\phi^0$ . Mass eigenstates of these  $U(1)_{\text{DM}}$ -charged scalars and their mixing angle  $\theta$  are obtained as

$$\begin{pmatrix} s_1^0 \\ s_2^0 \end{pmatrix} = \begin{pmatrix} \cos \theta & -\sin \theta \\ \sin \theta & \cos \theta \end{pmatrix} \begin{pmatrix} \eta^0 \\ s^0 \end{pmatrix}, \quad \sin 2\theta = \frac{\sqrt{2}\mu_3 v}{m_{s_2^0}^2 - m_{s_1^0}^2}. \quad (6)$$

Mass eigenvalues  $m_{s_1^0}$  and  $m_{s_2^0}$  of these neutral complex scalars are defined by

$$\begin{aligned} m_{s_1^0}^2 &= \frac{1}{2} \left( m_\eta^2 + m_s^2 - \sqrt{(m_\eta^2 - m_s^2)^2 + 2\mu_3^2 v^2} \right), \\ m_{s_2^0}^2 &= \frac{1}{2} \left( m_\eta^2 + m_s^2 + \sqrt{(m_\eta^2 - m_s^2)^2 + 2\mu_3^2 v^2} \right), \end{aligned} \quad (7)$$

where  $m_s^2 = \mu_s^2 + \lambda_{s\phi} v_\phi^2 / 2 + \lambda_{s\sigma} v_\sigma^2 / 2$  and  $m_\eta^2 = \mu_\eta^2 + (\lambda_{\phi\phi} + \lambda_{\eta\phi}) v_\phi^2 / 2 + \lambda_{\sigma\eta} v_\sigma^2 / 2$ . Finally, the mass of charged scalars  $\eta^\pm$  is

$$m_{\eta^\pm}^2 = \mu_\eta^2 + \lambda_{\phi\phi} \frac{v^2}{2} + \lambda_{\sigma\eta} \frac{v_\sigma^2}{2}. \quad (8)$$

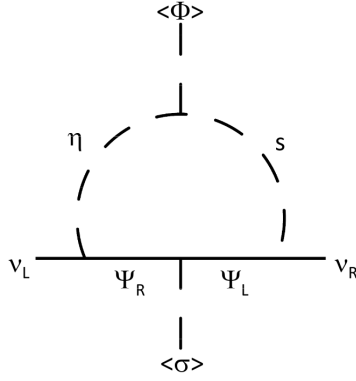


FIG. 1: Diagram for Dirac mass terms of neutrinos.

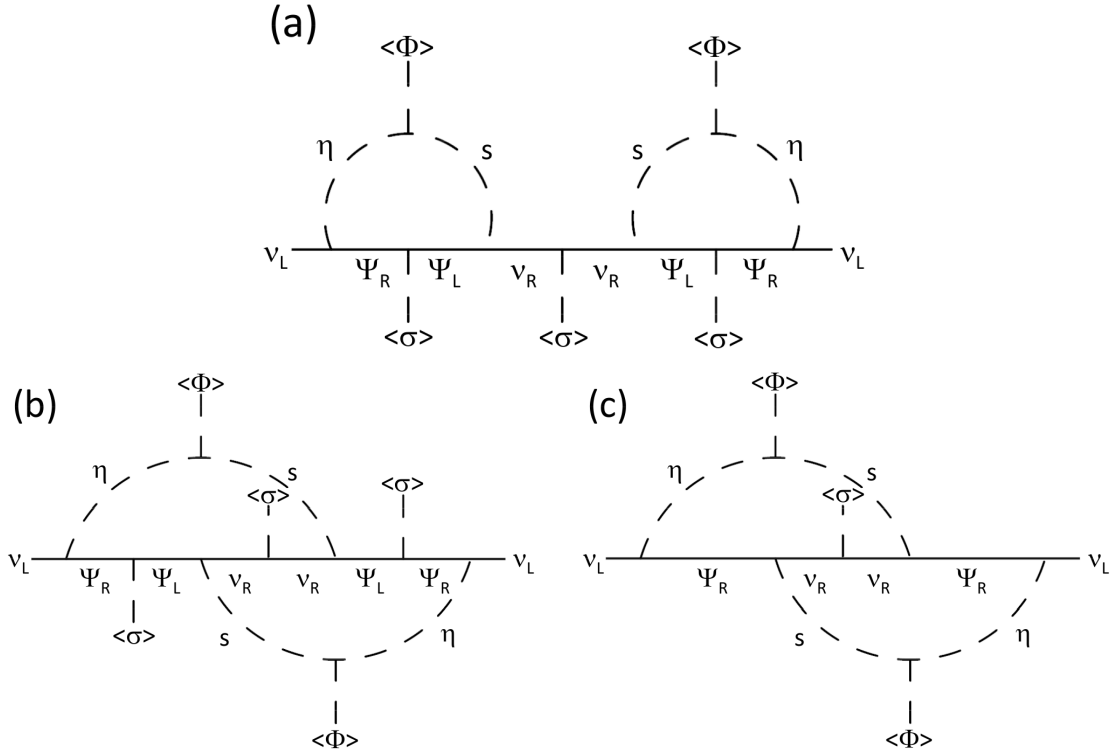


FIG. 2: Diagrams for light Majorana neutrino masses.

### III. NEUTRINO MASS AND DARK MATTER

#### A. Neutrino Mass

In this model, Dirac mass terms for neutrinos are generated by a one-loop diagram in Fig. 1. This diagram is used also in a model in ref. [14] in which lepton number is conserved.

Via the seesaw mechanism, tiny Majorana masses of light neutrinos are induced at the two-loop level as shown in Fig. 2(a) (See also refs. [11, 12]). In addition, there are one-particle-irreducible (1PI) diagrams also at the two-loop level (Figs. 2(b) and 2(c))<sup>2</sup>. The Majorana mass matrix is calculated as

$$(m_\nu)_{\ell\ell'} = \left(\frac{1}{16\pi^2}\right)^2 \left\{ \sum_{i,j,a} f_{\ell i} h_{ia} (m_R)_a (h^T)_{aj} (f^T)_{j\ell'} \left[ (I_1)_{ija} + \{I_2\}_{ija} \right] + \sum_{i,j,a} f_{\ell i} (y_3^\dagger)_{ia} (m_R)_a (y_3^*)_{aj} (f^T)_{j\ell'} \{I_3\}_{ija} \right\}, \quad (9)$$

where dimensionless functions  $I_1$ ,  $I_2$ , and  $I_3$  are defined by

$$(I_1)_{ija} = -\frac{(8\pi^2 \sin 2\theta)^2 m_{\Psi_i} m_{\Psi_j}}{(m_R)_a^2} \left[ \int \frac{d^4 k_1}{(2\pi)^4} \frac{1}{k_1^2 - m_{\Psi_i}^2} \left\{ \frac{1}{k_1^2 - m_{s_1^0}^2} - \frac{1}{k_1^2 - m_{s_2^0}^2} \right\} \right] \times \left[ \int \frac{d^4 k_2}{(2\pi)^4} \frac{1}{k_2^2 - m_{\Psi_j}^2} \left\{ \frac{1}{k_2^2 - m_{s_1^0}^2} - \frac{1}{k_2^2 - m_{s_2^0}^2} \right\} \right], \quad (10)$$

$$(I_2)_{ija} = (8\pi^2 \sin 2\theta)^2 m_{\Psi_i} m_{\Psi_j} \times \iint \frac{d^4 k_1}{(2\pi)^4} \frac{d^4 k_2}{(2\pi)^4} \left\{ \frac{1}{k_1^2 - m_{s_1^0}^2} - \frac{1}{k_1^2 - m_{s_2^0}^2} \right\} \frac{1}{k_1^2 - m_{\Psi_i}^2} \times \frac{1}{(k_1 - k_2)^2 - (m_R)_a^2} \left\{ \frac{1}{k_2^2 - m_{s_1^0}^2} - \frac{1}{k_2^2 - m_{s_2^0}^2} \right\} \frac{1}{k_2^2 - m_{\Psi_j}^2}, \quad (11)$$

$$(I_3)_{ija} = (8\pi^2 \sin 2\theta)^2 \iint \frac{d^4 k_1}{(2\pi)^4} \frac{d^4 k_2}{(2\pi)^4} k_1 \cdot k_2 \left\{ \frac{1}{k_1^2 - m_{s_1^0}^2} - \frac{1}{k_1^2 - m_{s_2^0}^2} \right\} \frac{1}{k_1^2 - m_{\Psi_i}^2} \times \frac{1}{(k_1 - k_2)^2 - (m_R)_a^2} \left\{ \frac{1}{k_2^2 - m_{s_1^0}^2} - \frac{1}{k_2^2 - m_{s_2^0}^2} \right\} \frac{1}{k_2^2 - m_{\Psi_j}^2}, \quad (12)$$

which correspond to the diagrams in Figs. 2(a), 2(b), and 2(c), respectively.

If there is only one  $\Psi$  (one  $\Psi_L$  and one  $\Psi_R$ ), the mass matrix  $(m_\nu)_{\ell\ell'}$  is proportional to  $f_\ell f_{\ell'}$ . Then two of three eigenvalues of  $(m_\nu)_{\ell\ell'}$  are zero and the mass matrix conflicts with the oscillation data. Therefore two  $\Psi_i$  are introduced in this model. We also introduce two  $(\nu_R)_i$  in order for an easy search of parameter sets which satisfy experimental constraints<sup>3</sup>. Then the rank of  $(m_\nu)_{\ell\ell'}$  is two, for which one neutrino becomes massless. Hereafter degeneracy of right-handed neutrino masses,  $m_R \equiv (m_R)_1 = (m_R)_2$ , is assumed for simplicity.

<sup>2</sup> Such 1PI diagrams were overlooked in refs. [11, 12].

<sup>3</sup> With two  $\Psi_i$ , one  $\nu_R$  is sufficient for that the rank of  $(m_\nu)_{\ell\ell'}$  is two.

The mass matrix  $(m_\nu)_{\ell\ell'}$  is diagonalized by the Maki-Nakagawa-Sakata (MNS) matrix  $U_{\text{MNS}}$  as  $U_{\text{MNS}}^T m_\nu U_{\text{MNS}} = \text{diag}(m_1, m_2, m_3)$ . The standard parametrization of the MNS matrix is

$$U_{\text{MNS}} = \begin{pmatrix} 1 & 0 & 0 \\ 0 & c_{23} & s_{23} \\ 0 & -s_{23} & c_{23} \end{pmatrix} \begin{pmatrix} c_{13} & 0 & s_{13} e^{-i\delta} \\ 0 & 1 & 0 \\ -s_{13} e^{i\delta} & 0 & c_{13} \end{pmatrix} \begin{pmatrix} c_{12} & s_{12} & 0 \\ -s_{12} & c_{12} & 0 \\ 0 & 0 & 1 \end{pmatrix}, \quad (13)$$

where  $c_{ij}$  and  $s_{ij}$  stand for  $\cos \theta_{ij}$  and  $\sin \theta_{ij}$ , respectively. Mixing angles  $\theta_{ij}$  and  $\Delta m_{ij}^2 \equiv m_i^2 - m_j^2$  are constrained by neutrino oscillation measurements [1–5]. In our analyses we use the following values as an example;

$$s_{23}^2 = \frac{1}{2}, \quad s_{13}^2 = 0, \quad s_{12}^2 = \frac{1}{3}, \quad (14)$$

$$\Delta m_{21}^2 = 7.5 \times 10^{-5} \text{ eV}^2, \quad |\Delta m_{31}^2| = 2.3 \times 10^{-3} \text{ eV}^2. \quad (15)$$

Notice that there is no difficulty to use nonzero values of  $s_{13}$  [32] in our analyses. In Table II, we show two examples (Set A and Set B) for the parameter set which reproduces the values given in eqs. (14) and (15) for  $\Delta m_{31}^2 > 0$ . These sets satisfy also other experimental constraints as shown below.

## B. Lepton Flavor Violation

Charged scalar bosons  $\eta^\pm$ , which also have a  $U(1)_{\text{DM}}$ -charge, contribute to the  $\mu \rightarrow e\gamma$  process. The branching ratio of  $\mu \rightarrow e\gamma$  is calculated as

$$\text{BR}(\mu \rightarrow e\gamma) = \frac{3\alpha_{\text{EM}}}{64\pi G_F^2} \left| \frac{1}{m_{\eta^\pm}^2} f_{\mu i} F_2 \left( \frac{m_{\Psi_i}^2}{m_{\eta^\pm}^2} \right) (f^\dagger)_{ie} \right|^2, \quad (16)$$

where

$$F_2(a) \equiv \frac{1 - 6a + 3a^2 + 2a^3 - 6a^2 \ln(a)}{6(1-a)^4}. \quad (17)$$

For Set A and Set B, we obtain  $\text{BR}(\mu \rightarrow e\gamma) = 5.1 \times 10^{-13}$  and  $1.7 \times 10^{-12}$ , respectively. They satisfy the current upper bound;  $\text{BR}(\mu \rightarrow e\gamma) < 2.4 \times 10^{-12}$  (90% CL) [33]. These values for Set A and Set B could be within the future experimental reach.



	Set A	Set B
$f_{\ell i}$	$\begin{pmatrix} -0.00726 & 0.00667 \\ -0.0523 & 0.0206 \\ -0.0378 & 0.00723 \end{pmatrix}$	$\begin{pmatrix} -0.0485 & 0.0505 \\ -0.0364 & 0.0433 \\ 0.0606 & -0.0577 \end{pmatrix}$
$h_{ij}$	$\begin{pmatrix} -0.119 & 0.150 \\ 0.150 & 0.150 \end{pmatrix}$	$\begin{pmatrix} 0.544 & 0.505 \\ 0.505 & 0.505 \end{pmatrix}$
$(y_3)_{ij}$	$\begin{pmatrix} 0.0152 & 0.0152 \\ 0.0152 & 0.0152 \end{pmatrix}$	$\begin{pmatrix} 0.0101 & 0.0101 \\ 0.0101 & 0.0101 \end{pmatrix}$
$m_R \equiv (m_R)_1 = (m_R)_2$	250 GeV	200 GeV
$\{m_{\Psi_1}, m_{\Psi_2}\}$	{57.0 GeV, 800 GeV}	{800 GeV, 800 GeV}
$\{m_{h^0}, m_{H^0}, \cos \alpha\}$	{120 GeV, 140 GeV, $1/\sqrt{2}$ }	{130 GeV, 300 GeV, 1}
$\{m_{s_1^0}, m_{s_2^0}, \cos \theta\}$	{200 GeV, 300 GeV, 0.05}	{55.0 GeV, 250 GeV, 0.05}
$m_{\eta^\pm}$	280 GeV	220 GeV
$g_{B-L}$	0.2	0.2
$m_{Z'}$	2000 GeV	2000 GeV

TABLE II: Two examples of the parameter set which satisfies experimental constraints. The dark matter is  $\Psi_1$  for Set A while it is  $s_1^0$  ( $\simeq s^0$ ) for Set B.

### C. Dark Matter

The dark matter candidate is  $\Psi_1$  (for Set A) or  $s_1^0$  (for Set B). The relic abundance of dark matter is constrained stringently by the WMAP experiment as  $\Omega_{\text{DM}} h^2 \simeq 0.11$  [16]. The dark matter candidate in this model pair-annihilates into a pair of SM fermions  $f$  by  $s$ -channel processes mediated by  $h^0$  and  $H^0$  for both of Set A and Set B. The  $t$ -channel diagram is highly suppressed due to the small values of  $f_{\ell i}$ , which are required by the  $\mu \rightarrow e\gamma$  search results.

We first consider the case where  $\Psi_1$  is the dark matter, i.e. Set A, whose mass is given by  $(y_\Psi)_1 v_\sigma / \sqrt{2}$ . Because  $v_\sigma > 3.5$  TeV, the magnitude of  $(y_\Psi)_1$  is  $\lesssim 0.01$  for  $m_{\Psi_1} = \mathcal{O}(10)$  GeV. The annihilation cross section is suppressed by the small  $(y_\Psi)_1$  because it is proportional to

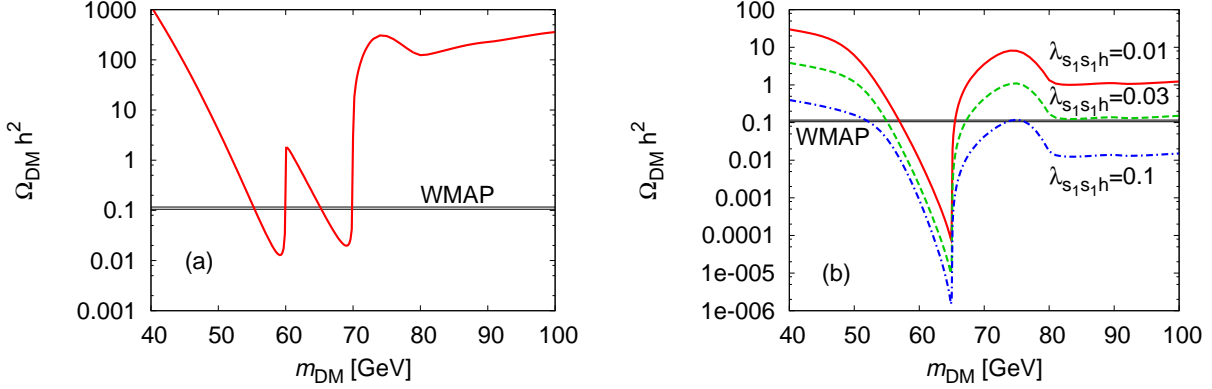


FIG. 3: The relic abundance with respect to the dark matter mass  $m_{\text{DM}}$ . Horizontal lines show  $1\sigma$  allowed region ( $0.1053 \leq \Omega_{\text{DM}} h^2 \leq 0.1165$ ) of the WMAP result. (a) Case for Set A where  $\Psi_1$  is the dark matter candidate. (b) Case for Set B where  $s_1^0$  is the dark matter candidate. Three curves are obtained for  $\lambda_{s_1 s_1 h} = 0.1$  (lower curve), 0.03 (middle curve), and 0.01 (upper curve).

$(m_f(y_\Psi)_1 \sin 2\alpha/v)^2$ . In order to enhance the cross section for the appropriate relic abundance, a large mixing between  $\sigma_r^0$  (which couples with  $\Psi_1$ ) and  $\phi_r^0$  (which couples with SM fermion  $f$ ) is required [23, 24]. Thus we take  $\cos \alpha = 1/\sqrt{2}$  for Set A. Then the WMAP result gives a constraint on the dark matter mass  $m_{\Psi_1}$  for fixed  $m_{h^0}$  and  $m_{H^0}$  (120 GeV and 140 GeV for Set A, respectively). Figure 3(a) shows dependence of the relic abundance on  $m_{\Psi_1}$  where other parameters are the same as values of Set A. It can be confirmed in the figure that our choice  $m_{\Psi_1} = 57$  GeV is consistent with the WMAP result for Set A.

Next, we consider the case where  $s_1^0$  is the dark matter. A coupling constant  $\lambda_{s_1 s_1 h}$  for  $\lambda_{s_1 s_1 h} v s_1^0 (s_1^0)^* h^0$  is constrained by the WMAP result. Figure 3(b) shows the relic abundance of  $s_1^0$  as functions of  $m_{s_1^0}$  for several values of  $\lambda_{s_1 s_1 h}$ . In the figure, we used  $m_{h^0} = 130$  GeV and  $m_{H^0} = 300$  GeV which are the values for Set B. Contribution of  $H^0$  to the annihilation cross section is negligible because  $m_{H^0}$  is taken to be away from  $2m_{s_1^0} = 110$  GeV for simplicity. In order to satisfy the WMAP constraint for Set B, we find that

$$\lambda_{s_1 s_1 h} \simeq 0.03. \quad (18)$$

This constraint can be satisfied easily because  $\lambda_{s_1 s_1 h}$  can be taken to be arbitrary depending on several parameters in the scalar potential.

Finally, we discuss the constraint from direct search experiments for the dark matter. If  $s_1^0$  is mainly composed of  $\eta^0$ , it cannot be a viable candidate for the dark matter even if it

is the lightest  $U(1)_{\text{DM}}$ -charged particle. This is because the scattering cross section with a nucleon  $N$  ( $N = p, n$ ) becomes too large due to the weak interaction. Thus  $s_1^0$  should be dominantly made from singlet  $s^0$ , and this is the reason why we take  $\cos \theta = 0.05$  for Set B. Scattering cross sections of dark matter candidates ( $\Psi_1$  and  $s_1^0$ ) with a nucleon  $N$  ( $N = p, n$ ) are given by

$$\sigma(\Psi_1 N \rightarrow \Psi_1 N) \simeq \frac{8g_{\text{B-L}}^2 m_{\Psi_1}^2 \sin^2 \alpha \cos^2 \alpha}{v^2 m_{Z'}^2} \left( \frac{1}{m_{h^0}^2} - \frac{1}{m_{H^0}^2} \right)^2 \frac{m_N^2 m_{\Psi_1}^2}{\pi (m_{\Psi_1} + m_N)^2} f_N^2 + \left( \frac{g_{\text{B-L}}}{m_{Z'}} \right)^4 \frac{m_{\Psi_1}^2 m_N^2}{4\pi (m_{\Psi_1} + m_N)^2}, \quad (19)$$

$$\sigma(s_1^0 N \rightarrow s_1^0 N) \simeq \frac{1}{4} \left\{ \frac{\lambda_{s_1 s_1 h} \cos \alpha}{m_{h^0}^2} + \frac{\lambda_{s_1 s_1 H} \sin \alpha}{m_{H^0}^2} \right\}^2 \frac{m_N^2}{\pi (m_{s_1^0} + m_N)^2} f_N^2 + \left( \frac{g_{\text{B-L}}}{m_{Z'}} \right)^4 \frac{m_{s_1^0}^2 m_N^2}{4\pi (m_{s_1^0} + m_N)^2}, \quad (20)$$

$$f_N \equiv \sum_{q=u,d,s} m_N f_{Tq} + \frac{2}{9} m_N f_{Tg}, \quad (21)$$

where  $m_N$  is the mass of the nucleon and we use  $f_{Tu} + f_{Td} = 0.056$ ,  $f_{Ts} = 0$  [34], and  $f_{Tg} = 0.944$  [35]. Our results for the  $Z'$  mediation are consistent with those in ref. [36]. Difference between  $p$  and  $n$  is neglected. We have  $\sigma(\Psi_1 N \rightarrow \Psi_1 N) = 2.7 \times 10^{-45} [\text{cm}^2]$  for  $m_{\Psi_1} = 57 \text{ GeV}$  for Set A. The value is dominantly given by the  $Z'$  mediation while scalar mediations give only  $2.4 \times 10^{-48} [\text{cm}^2]$ . On the other hand, for  $m_{s_1^0} = 55 \text{ GeV}$  for Set B, we have  $\sigma(s_1^0 N \rightarrow s_1^0 N) = 4.4 \times 10^{-45} [\text{cm}^2]$  by taking into account eq. (18) as the WMAP constraint. Contributions from  $Z'$  and  $h^0$  are  $2.6 \times 10^{-45} [\text{cm}^2]$  and  $1.7 \times 10^{-45} [\text{cm}^2]$ , respectively. These values of cross sections for two sets are just below the constraint from the XENON100 experiment [37]. Notice that even if such values are excluded in near future this model is not ruled out because the  $Z'$  contribution ( $\propto v_\sigma^{-4}$ ) can be easily suppressed by a little bit larger  $v_\sigma$ .

## IV. PROSPECTS FOR COLLIDER PHENOMENOLOGY AND DISCUSSION

### A. Collider Phenomenology

We have seen that Set A and Set B satisfy experimental constraints in the previous section. Let us discuss the collider phenomenology by using these parameter sets.

	BR( $Z' \rightarrow XX$ )								
	$q\bar{q}$	$\ell^-\ell^+$	$\nu_L\bar{\nu}_L$	$\nu_R\bar{\nu}_R$	$\Psi_1\bar{\Psi}_1$	$\Psi_2\bar{\Psi}_2$	$s_1^0(s_1^0)^*$	$s_2^0(s_2^0)^*$	$\eta^+\eta^-$
Set A	0.20	0.30	0.15	0.10	0.13	0.085	0.012	0.011	0.011
Set B	0.21	0.31	0.16	0.10	0.089	0.089	0.013	0.012	0.012

TABLE III: Branching ratios of  $Z'$  decays.

Since  $U(1)_{B-L}$  is dealt with as the origin of neutrino masses etc.,  $Z'$  is an important particle in this model. The  $Z'$  can be the mother particle to produce the  $U(1)_{DM}$ -charged particles and  $\nu_R$  in the model at collider experiments because they are all  $U(1)_{B-L}$ -charged. For  $m_{Z'} = 2 \text{ TeV}$  and  $g_{B-L} = 0.2$ , the production cross section of  $Z'$  at the CERN Large Hadron Collider (LHC) with  $\sqrt{s} = 14 \text{ TeV}$  is  $70 \text{ fb}$  [25, 28]. The number of  $Z'$  produced with  $100 \text{ fb}^{-1}$  is 7000. Branching ratios of the  $Z'$  decay are shown in Table III for Set A and Set B. Similar  $\text{BR}(Z' \rightarrow XX)$  are predicted for the two sets because  $Z'$  is sufficiently heavier than the others. The large  $\text{BR}(Z' \rightarrow \ell\bar{\ell})$  (cf.  $\text{BR}(Z \rightarrow \ell\bar{\ell}) \simeq 0.1$  in the SM) could be utilized for discovery of the  $Z'$  at the LHC. The B–L nature of the  $Z'$  would be tested if  $\text{BR}(Z' \rightarrow b\bar{b})/\text{BR}(Z' \rightarrow \mu\bar{\mu}) = 1/3$  would be confirmed.

Since each  $U(1)_{DM}$ -charged particle decays finally into a dark matter, 25 % of the  $Z'$  gives a pair of the dark matter ( $\Psi_1\bar{\Psi}_1$ ) for Set A while 22 % of the  $Z'$  produces  $s_1^0(s_1^0)^*$  for Set B. Since  $f_{\ell i}$  are preferred to be small in order to satisfy the constraint on  $\text{BR}(\mu \rightarrow e\gamma)$ ,  $\Psi_2$  for Set A ( $\Psi_1$  and  $\Psi_2$  for Set B) decays into  $\nu_R s_1^0$  with  $h_{ij}$  and  $(y_3)_{ij}$ . Subsequently  $\nu_R$  dominantly decays into  $W^\pm\ell^\mp$  or  $Z\nu_L$  (See Table IV), and thus it gives a visible signal. For Set A,  $s_2^0 (\simeq \eta^0)$  decays invisibly into  $\Psi_1\bar{\nu}_L$  with  $f_{ij}$ , and the  $s_1^0 (\simeq s^0)$  decays also invisibly into  $\Psi_1\bar{\nu}_L$  with  $f_{ij} \cos\theta$ . For Set B,  $s_2^0$  decays into  $s_1^0 h^0$  with  $\mu_3$ . It is clear that  $\eta^\pm$  provide visible signals with  $\eta^\pm \rightarrow \ell^\pm\bar{\Psi}_1$  for Set A and  $W^\pm s_1^0$  for Set B. As a result, about 30 % (36 %) of the  $Z'$  for Set A (B) is invisible and the constraint on  $m_{Z'}$  becomes milder. If a light  $Z'$  ( $\lesssim 1 \text{ TeV}$ ) is discovered at the LHC by  $Z' \rightarrow \ell\bar{\ell}$ , this model could be tested further at future linear colliders by measuring the amount of invisible decay of the  $Z'$ .

It is a good feature of this model that a light  $\nu_R$  is acceptable naturally because of loop-suppressed Dirac mass terms. For Set A, we see that  $Z' \rightarrow \nu_R\bar{\nu}_R$  and  $Z' \rightarrow \Psi_2\bar{\Psi}_2$  followed by  $\Psi_2 \rightarrow s_2^0\nu_R$  ( $s_2^0\bar{\nu}_R$ ) produce about 1200 pairs of  $\nu_R$  from 7000 of  $Z'$ . For Set B, the number of  $\nu_R$  pairs increases to about 1700 because of an additional contribution from the decay

	BR( $\nu_R \rightarrow XY$ )			
	$W^\pm \ell^\mp$	$Z \nu_L$	$h^0 \nu_L$	$H^0 \nu_L$
Set A	0.53	0.28	0.10	0.09
Set B	0.52	0.28	0.21	0

TABLE IV: Branching ratios of  $\nu_R$  decays for Set A and Set B where  $(m_R)_1 = (m_R)_2$  is assumed. For Set B, the decay  $\nu_R \rightarrow H^0 \nu_L$  is kinematically forbidden.

of  $\Psi_1$ . The  $\nu_R$  decays into  $W^\pm \ell^\mp$ ,  $Z \nu_L$ ,  $h^0 \nu_L$ , and  $H^0 \nu_L$  through mixing due to the 1-loop induced Dirac Yukawa coupling. For Set A and Set B,  $\nu_R \rightarrow W^\pm \ell^\mp$  is the main decay mode as shown in Table IV. Then, the Majorana mass of  $\nu_R$  can be reconstructed by observing the invariant mass of the  $jj\ell^\pm$  [38, 39]. The  $h^0$  produced from  $\nu_R$  will be energetic due to the helicity structure. Therefore it would be possible to test the existence of  $\nu_R$  by search for energetic  $b\bar{b}$ .

If  $\Psi_1$  is the dark matter, the Yukawa coupling constant  $y_{\Psi_1 \Psi_1 h} [= (y_\Psi)_1 \sin \alpha]$  for the decay  $h^0 \rightarrow \Psi_1 \bar{\Psi}_1$  should be small because of small  $m_{\Psi_1}$ . For Set A, we have  $y_{\Psi_1 \Psi_1 h} \simeq 0.01$ . Then main decay mode of  $h^0$  ( $m_{h^0} < 2m_t$ ) is  $h^0 \rightarrow b\bar{b}$  similarly to the SM case. Since  $\sin \alpha = \mathcal{O}(1)$  is required to obtain the appropriate relic abundance of  $\Psi_1$ , the main decay mode of  $H^0$  is also  $H^0 \rightarrow b\bar{b}$ . Their decay widths are about a half of the width of the SM Higgs boson. The large mixing prefers that  $m_{h^0}$  and  $m_{H^0}$  are of the same order of magnitude. Thus we would find two SM-like Higgs bosons whose masses are  $\mathcal{O}(100)$  GeV, e.g.  $m_{h^0} = 120$  GeV and  $m_{H^0} = 140$  GeV for Set A.

On the other hand, if  $s_1^0$  is the dark matter, the interaction of  $h^0$  with dark matter  $s_1^0$  should satisfy eq. (18). Then the invisible decay  $h^0 \rightarrow s_1^0 (s_1^0)^*$  dominates for  $2m_{s_1^0} \leq m_{h^0} < 2m_W$ . We have  $\text{BR}(h^0 \rightarrow s_1^0 (s_1^0)^*) = 0.38$  for Set B where  $h^0$  is 100 % SM-like. Therefore, even if  $h^0$  is not discovered at the LHC within a year, a light  $h^0$  might exist because the constraint on  $m_{h^0}$  is relaxed from the one for the SM Higgs boson. If such a light  $h^0$  is discovered late with smaller number of signals than the SM expectation, this model would be confirmed at the LHC [40, 41] and future linear colliders [42] by “observing” the  $h^0$  invisible decay<sup>4</sup>.

<sup>4</sup> When  $m_{h^0} < 2m_{s_1^0}$ , the Higgs portal dark matter such as the  $s_1^0$  in this scenario would be able to be

## B. Some remarks

In our paper, we have assumed the mass of the  $Z'$  boson to be 2 TeV. It is expected that the lower bound on the mass will be more and more stringent due to the new results from the LHC and that the above value of the mass would be excluded in near future. In such a case, the mass should be taken to be higher values than 2 TeV. Then, the branching ratio of  $Z' \rightarrow \Psi_2 \bar{\Psi}_2$  becomes closer values to that of  $Z' \rightarrow \Psi_1 \bar{\Psi}_1$  because the effect of their mass difference becomes less significant. Heavy  $Z'$  is achieved by assuming larger values of  $g_{B-L}$  or assuming larger values of  $v_\sigma$ . For the former case, most of the phenomenological analyses are unchanged. For the latter case, our phenomenological analysis would be changed slightly. Even in this case, the experimental constraints from neutrino experiments and the  $\mu \rightarrow e\gamma$  results can be satisfied by using smaller values of  $(y_R)_i$  and  $(y_\Psi)_i$  which keep  $(m_R)_i$  and  $m_{\Psi_i}$  unchanged from values in our Sets A and B. In the scenario of Set A, a smaller value of  $(y_\Psi)_1$  results in a larger value of the DM abundance which is proportional to  $(y_\Psi)_1^{-2}$ . Even if the red curve in Fig. 3(a) goes up by about a factor of 10, the WMAP result can still be explained by  $m_{\Psi_1} \simeq 60$  GeV. Therefore, we can accept three times larger  $v_\sigma$  (namely, three times smaller  $(y_\Psi)_1$ ) than the value (5 TeV) we used.

A shortage is about the gauge anomaly for  $U(1)_{B-L}$ . It is well known that  $U(1)_{B-L}$  is free from anomaly if three singlet fermions of  $B-L = -1$  (usual right-handed neutrinos) are added to the SM. However, we have introduced not three but only two  $(\nu_R)_i$ , and their  $B-L$  is not  $-1$  but  $+1$ . There are also extra  $U(1)_{B-L}$ -charged fermions  $((\Psi_L)_i$  and  $(\Psi_R)_i$ ). Therefore the  $U(1)_{B-L}$  gauge symmetry has anomalies for the triangle diagrams of  $[U(1)_{B-L}]^3$  and  $[U(1)_{B-L}] \times [\text{gravity}]^2$ . They would be resolved by some heavy singlet fermions with appropriate  $B-L$  (See e.g., ref. [45]).

Another is the way to reproduce the baryon asymmetry of the universe. Since we have used only particles below the TeV-scale, leptogenesis [46] does not work in a natural way. Heavy fermions to eliminate the  $U(1)_{B-L}$  gauge anomaly might help. The electroweak baryogenesis [47] would be accommodated by the introduction of an additional Higgs doublet to the model so that new source of CP-violation appears in the Higgs potential.

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tested at the LHC [43] and future linear colliders [44].

## V. CONCLUSIONS

We have presented the TeV-scale seesaw model in which  $U(1)_{B-L}$  gauge symmetry can be the common origin of neutrino masses, the dark matter mass (if  $\Psi_1$  is the one), and stability of the dark matter. Light neutrino masses are generated by the two-loop diagrams which are also contributed by the dark matter, a Dirac fermion  $\Psi_1$  or a complex scalar  $s_1^0$ . The symmetry to stabilize the dark matter appears in the  $U(1)_{B-L}$ -symmetric Lagrangian without introducing additional global symmetry (ex. a  $Z_2$  symmetry). It has been shown that this model can be compatible with constraints from the neutrino oscillation data, the search for  $\mu \rightarrow e\gamma$ , the relic abundance of the dark matter, and the direct search results for dark matter. It should be emphasized that these constraints are satisfied with sizable coupling constants ( $\gtrsim 10^{-2}$ ) and new particles (including  $\nu_R$ ) whose masses are at or below the TeV-scale.

We have discussed collider phenomenology in this model by using two sets of parameters which satisfy constraints from the current experimental data. The  $U(1)_{B-L}$  gauge boson  $Z'$  can be discovered at the LHC by observing  $Z' \rightarrow \ell\bar{\ell}$  if it is not too heavy. In our model, since the dark matter has  $U(1)_{B-L}$  charge,  $Z'$  partially decays into a pair of the dark matter. As a result, more than 30 % of produced  $Z'$  is invisible for both the sets. Then a lighter  $Z'$  is allowed than the usual experimental bound. Detailed studies for such a  $Z'$  would be performed at future linear colliders.

Since masses of  $(\nu_R)_i$  and  $\Psi_i$  are obtained by the  $U(1)_{B-L}$  breaking, it would be natural that  $\Psi_1$  is light and becomes the dark matter when we assume  $\nu_R$  masses are of  $\mathcal{O}(100)$  GeV. In this case, a large mixing between  $h^0$  and  $H^0$  is required for the appropriate relic abundance because the Yukawa coupling constant is small. Decay branching ratios of  $h^0$  and  $H^0$  are almost the same as the one for the SM Higgs boson. Therefore two SM-like Higgs bosons with similar masses (ex.  $m_{h^0} = 120$  GeV and  $m_{H^0} = 140$  GeV for Set A) would be discovered at the LHC.

On the other hand, if  $s_1^0$  is the dark matter, the decay of the lighter Higgs boson  $h^0$  can be dominated by invisible  $h^0 \rightarrow s_1^0(s_1^0)^*$ . For Set B, we obtain  $\text{BR}(h^0 \rightarrow s_1^0(s_1^0)^*) = 38\%$ . Then this model would be tested at future linear colliders by measuring the amount of the invisible decay.

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